Regular particle acceleration in relativistic jets *

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Abstract.
Exact solution is obtained for electromagnetic field around a conducting cylinder of infinite length and finite radius, with a periodical axial current, when the wave length is much larger than the radius of the cylinder. The solution describes simultaneously the fields in the near zone close to the cylinder, and transition to the wave zone. Proper long-wave oscillations of such cylinder are studied. The electromagnetic energy flux from the cylinder is calculated. These solutions could be applied for description of the electromagnetic field around relativistic jets from active galactic nuclei and quasars and particle acceleration inside jets.

Keywords: accretion disk, X-ray source, jet

1. Introduction

Objects of different scale and nature in the universe: from young and very old stars to active galactic nuclei (AGN) (Bridle, 1984),(Eilek, 1984),(Eilek and Hughes, 1990) show existence of collimated outbursts - jets. Geometrical sizes of jets lay between parsecs and megaparsecs. The origin of jets is not well understood and only several qualitative mechanisms are proposed. Theory of jets should answer to the question of the origin of relativistic particles in outbursts from AGN, where synchrotron emission is observed. Relativistic particles, ejected from the central machine rapidly loose their energy so the problem arises of particle acceleration inside the jet, see reviews (Begelman et al., 1984),(Bisnovatyi-Kogan, 1993).

It is convenient sometimes to investigate jets in a simple model of infinitely long circular cylinder (Chandrasekhar and Fermi, 1953). Magnetic fields in collimated jets is determining its direction, and axial current is stabilizing its elongated form at large distances from the source (AGN) (Bisnovatyi-Kogan et al., 1969). When observed with high angular resolution these jets show structure with bright knots separated by relatively dark regions (Bridle, 1984), (Eilek, 1984), citetmw. High percentages of polarization, sometimes more then 50% in some objects, indicates the nonthermal nature of the radiation which is well

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explained as the synchrotron radiation of the relativistic electrons in a weak but ordered magnetic field. Estimation of the life time of these electrons, based on the observed luminosities and spectra, often gives values much less the their kinematic ages t_k/c, where d is the distance of the emitting point from the central source. Because the jet flow most likely originated from an outburst or continuous outflow from the central source, there is a necessity of continuous reacceleration of electrons in the jets in order to explain the observations. Acceleration mechanism for electrons in extragalactic jets, proposed in (Bisnovatyi-Kogan and Lovelace, 1995) considers that intense long-wavelength electromagnetic oscillations accompany a relativistic jet as a result of the non-steady mechanism of the jet's generation in the nucleus of the source. The electromagnetic wave amplitudes envisioned are sufficient to give in situ acceleration of electrons to the very high energies observed $> 10^{13}$ eV. It was assumed that jets are formed by a sequence of outbursts from the nucleus with considerable charge separation at the moment of the outburst (Bisnovatyi-Kogan et al., 1969). The direction of motion the outburst is determined by the large-scale magnetic field. The outburst are accompanied by an intense electromagnetic disturbance which propagates outward moving with the jet material in the direction of the large scale magnetic field. It was suggested in (Bisnovatyi-Kogan et al., 1969) that a toroidal magnetic field, generated during the outbursts is important for the lateral confinement of of the jet.

When the plasma density in surrounding medium is small, the electromagnetic wave generated by the nonpotential plasma oscillations of the confined body is emitted outside and can accelerate particles. When the emitted wave is strong enough it washed out the medium around and the density can become very small, consisting only of the accelerated particles. The action of the oscillating knot is similar to the action of the pulsar, as inclined magnetic rotator. Both emit strong electromagnetic waves, which could effectively accelerate particles (Pacini, 1967), (Gunn and Ostriker, 1970). Long-periodic proper oscillations in the plasma cylinder with a finite radius, and emission of electromagnetic waves had been studied in (Bisnovatyi-Kogan and Lovelace, 1995), and in a simpler model in (Bisnovatyi-Kogan, 1996). Enhanced oscillations in such cylinder have been studied in (Bisnovatyi-Kogan, 2004). Both models are represented below.

2. Cylinder with oscillating current

Consider infinitely conducting circular cylinder in vacuum. This model is valid at low density in surrounding plasma, which cannot screen the

emitting electromagnetic wave. Maxwell equations are (Landau and Lifshits, 1982)

$$\mathrm{div}\,\mathbf{B} = 0, \quad \mathrm{rot}\,\mathbf{B} = \frac{1}{c}\frac{\partial \mathbf{E}}{\partial t} + \frac{4\pi}{c}\mathbf{j}, \tag{1}$$

$$\mathrm{rot}\,\mathbf{E} = -\frac{1}{c}\frac{\partial \mathbf{B}}{\partial t}, \quad \mathrm{div}\,\mathbf{E} = 4\pi\rho_e. \tag{2}$$

For periodic oscillations with all values $\sim \exp(-i\omega t)$ they read

$$\mathrm{div}\,\mathbf{B} = 0,\ \ \mathrm{rot}\,\mathbf{B} = -\frac{i\omega}{c}\mathbf{E} + \frac{4\pi}{c}\mathbf{j},\ \ \mathrm{rot}\,\mathbf{E} = \frac{i\omega}{c}\mathbf{B},\ \ \mathrm{div}\,\mathbf{E} = 0. \tag{3}$$

We use the same definitions for all complex values depending on coordinates. Consider infinitely long cylinder with zero charge density, where in the cylinder coordinate system (r, ϕ, z) the only nonzero components are $E_z,\ B_\phi,\ j_z$, and $\partial/\partial\phi = \partial/\partial z = 0$. Only two valid equations remain from the system (3):

$$\frac{dE_z}{dr} + \frac{i\omega}{c}B_\phi = 0, \quad \frac{1}{r}\frac{d(rB_\phi)}{dr} + \frac{i\omega}{c}E_z - \frac{4\pi}{c}j_z = 0; \tag{4}$$

from which, we obtain equation for E_z:

$$\frac{1}{r}\frac{d}{dr}\left(r\frac{dE_z}{dr}\right) + \frac{\omega^2}{c^2}E_z + \frac{4\pi i\omega}{c^2}j_z = 0. \tag{5}$$

3. Vacuum solution

In vacuum $j_z = 0$. Using non-dimensional variable $x = r\omega/c$, we obtain from (4),(5)

$$x^2 E_z'' + xE_z' + x^2 E_z = 0, \qquad B_\phi = iE_z'. \tag{6}$$

Here $'$ denote differentiation over x. The equation (6) belongs to Bessel type and has a solution

$$E_z = C_1 J_0(x) + C_2 Y_0(x), \qquad B_\phi = -i[C_1 J_1(x) + C_2 Y_1(x)]. \tag{7}$$

Relations for Bessel functions are (Gradshtein and Ryzhik, 1964)

$$J_0'(x) = -J_1(x), \qquad Y_0'(x) = -Y_1(x). \tag{8}$$

General solution for physical values, with account of the time dependence, is obtained from the real part of the complex solution at

$$\exp{(-i\omega t)} = \cos\omega t - i\sin\omega t, \qquad C_1 = C_1^{(r)} + iC_1^{(i)}, \quad C_2 = C_2^{(r)} + iC_i^{(i)}. \tag{9}$$

The general solution in vacuum is

$$E_z = [C_1^{(r)} J_0(x) + C_2^{(r)} Y_0(x)] \cos\omega t + [C_1^{(i)} J_0(x) + C_2^{(i)} Y_0(x)] \sin\omega t, \tag{10}$$

$$B_\phi = -[C_1^{(r)} J_1(x) + C_2^{(r)} Y_1(x)] \sin\omega t + [C_1^{(i)} J_1(x) + C_2^{(i)} Y_1(x)] \cos\omega t. \tag{11}$$

The boundary condition far from the cylinder follows from the demand that there exist only expanding wave. It means, that only functions depending on the combination $(x - \omega t)$ survive. Using asymptotic of Bessel functions at large arguments (Gradshtein and Ryzhik, 1964)

$$J_0(x) \approx \sqrt{\frac{2}{\pi x}} \cos(x - \frac{\pi}{4}), \quad J_1(x) \approx \sqrt{\frac{2}{\pi x}} \sin(x - \frac{\pi}{4}),$$

$$Y_0(x) \approx \sqrt{\frac{2}{\pi x}} \sin(x - \frac{\pi}{4}), \quad Y_1(x) \approx -\sqrt{\frac{2}{\pi x}} \cos(x - \frac{\pi}{4}), \quad \text{at} \quad x \gg 1. \tag{12}$$

we obtain for the expanding wave $C_1^{(i)} = -C_2^{(r)}$, $C_2^{(i)} = C_1^{(r)}$, leading to the following solution at large distances

$$E_z \approx \sqrt{\frac{2}{\pi x}} [C_1^{(r)} \cos(x - \frac{\pi}{4} - \omega t) + C_2^{(r)} \sin(x - \frac{\pi}{4} - \omega t)], \qquad B_\phi = -E_z. \tag{13}$$

The general vacuum solution, satisfying conditions at infinity reads as

$$E_z = [C_1^{(r)} J_0(x) + C_2^{(r)} Y_0(x)] \cos\omega t + [-C_2^{(r)} J_0(x) + C_1^{(r)} Y_0(x)] \sin\omega t, \tag{14}$$

$$B_\phi = -[C_1^{(r)} J_1(x) + C_2^{(r)} Y_1(x)] \sin\omega t + [-C_2^{(r)} J_1(x) + C_1^{(r)} Y_1(x)] \cos\omega t. \tag{15}$$

4. Solution inside the cylinder

The equation in the matter are

$$x^2 E_z'' + x E_z' + x^2 E_z + \frac{4\pi i}{\omega} x^2 j_z = 0, \qquad B_\phi = i E_z'. \tag{16}$$

A solution of the non-uniform linear equation (16) is a sum of a general solution of the uniform equation, and a particular solution of the non-uniform one $\mathcal{E}_0(x)$.

$$E_z = \mathcal{E}_1 J_0(x) + \mathcal{E}_2 Y_0(x) + \mathcal{E}_0(x). \tag{17}$$

The function $Y_0(x)$ is singular at $x = 0$, so for a finite solution $\mathcal{E}_2 = 0$. We look for a particular solution in the form $E_z = \mathcal{E}(x) J_0(x)$. First order equation with respect to $F = \mathcal{E}'$ follows from (16)

$$x^2 (F' J_0 + 2F J_0') + x F J_0 + \frac{4\pi i}{\omega} x^2 j_z = 0. \tag{18}$$

From this equation we obtain the general solution for the amplitude of the electric field in the matter, in presence of periodic EEF:

$$E_z = -\frac{4\pi i}{\omega} J_0(x) \int_0^x \frac{dy}{y J_0^2(y)} \int_0^y z J_0(z) j_z(z) dz + \mathcal{E}_1 J_0(x). \tag{19}$$

Consider waves much longer than the radius of the cylinder r_0

$$x_0 = \frac{\omega r_0}{c} \ll 1. \tag{20}$$

Than use expansion at $x \ll 1$ (Gradshtein and Ryzhik, 1964),

$$J_0 \approx 1 - \frac{x^2}{4}, \quad J_1 \approx \frac{x}{2}, \quad Y_0 \approx \frac{2}{\pi} \ln \frac{x}{2}, \quad Y_1 \approx -\frac{2}{\pi x}. \tag{21}$$

Using (21) we obtain from (19) the solution for long waves

$$E_z = -\frac{2i\omega}{c^2} \int_0^x I_z(y) \frac{dy}{y} + \mathcal{E}_1, \tag{22}$$

where $I_z(r) \equiv I_z(y)$ is the complex amplitude of the electrical current inside a cylindrical radius $r = cy/\omega$

$$I_z = 2\pi \int_0^r j_z r\, dr = 2\pi \frac{c^2}{\omega^2} \int_0^x j_z x\, dx \tag{23}$$

Complex values: the function $I_z(r)$ and the constant $\mathcal{E}_1$ are

$$I_z = I_z^{(r)} + i I_z^{(i)}, \qquad \mathcal{E}_1 = \mathcal{E}_1^{(r)} + i \mathcal{E}_1^{(i)}. \tag{24}$$

5. Matching of solutions and long-wave limit

The total electrical current through the cylinder $I_0 = I_z(r_0)$, and fields on its surface (inside) $E_0 = E_z(r_0)$, $B_0 = B_\phi(r_0)$ (real parts of complex relations) are defined as

$$I_0 = I_0^{(r)} \cos\omega t + I_0^{(i)} \sin\omega t, \quad E_0 =$$
$$= \left[\frac{2\omega}{c^2}\int_0^{x_0} I_z^{(i)}(y)\frac{dy}{y} + \mathcal{E}_1^{(r)}\right]\cos\omega t + \left[-\frac{2\omega}{c^2}\int_0^{x_0} I_z^{(r)}(y)\frac{dy}{y} + \mathcal{E}_1^{(i)}\right]\sin\omega t, \tag{25}$$
$$B_0 = \left[\frac{2\omega}{c^2}\frac{I_0^{(r)}}{x_0} + \frac{x_0}{2}\mathcal{E}_1^{(i)}\right]\cos\omega t + \left[\frac{2\omega}{c^2}\frac{I_0^{(i)}}{x_0} - \frac{x_0}{2}\mathcal{E}_1^{(r)}\right]\sin\omega t.$$

At small x_0 the external solution on the surface of the cylinder is written as

$$E_{z0} = [C_1^{(r)} + C_2^{(r)}\frac{2}{\pi}\ln\frac{x_0}{2}]\cos\omega t + [-C_2^{(r)} + C_1^{(r)}\frac{2}{\pi}\ln\frac{x_0}{2}]\sin\omega t, \tag{26}$$

$$B_{\phi 0} = -[C_1^{(r)}\frac{x_0}{2} - C_2^{(r)}\frac{2}{\pi x_0}]\sin\omega t + [-C_2^{(r)}\frac{x_0}{2} - C_1^{(r)}\frac{2}{\pi x_0}]\cos\omega t.$$

All field components are continuous at the cylinder surface in absence of the surface charges and currents. Matching magnetic and electrical fields we obtain using (25), (26) coefficients in the solution of the external electromagnetic field, which are determined by the periodic electrical current in the cylinder:

$$C_1^{(r)}\left(\frac{2}{\pi x_0} + \frac{x_0}{\pi}\ln\frac{x_0}{2}\right) = -\frac{2\omega}{c^2}\frac{I_0^{(r)}}{x_0} - \frac{x_0\omega}{c^2}\int_0^{x_0} I_z^{(r)}(y)\frac{dy}{y}, \tag{27}$$

$$C_2^{(r)}\left(\frac{2}{\pi x_0} + \frac{x_0}{\pi}\ln\frac{x_0}{2}\right) = \frac{2\omega}{c^2}\frac{I_0^{(i)}}{x_0} + \frac{x_0\omega}{c^2}\int_0^{x_0} I_z^{(i)}(y)\frac{dy}{y}. \tag{28}$$

$$\mathcal{E}_1^{(r)} = C_1^{(r)} + C_2^{(r)}\frac{2}{\pi}\ln\frac{x_0}{2} - \frac{2\omega}{c^2}\int_0^{x_0} I_z^{(i)}(y)\frac{dy}{y}, \tag{29}$$

$$\mathcal{E}_1^{(i)} = -C_2^{(r)} + C_1^{(r)}\frac{2}{\pi}\ln\frac{x_0}{2} + \frac{2\omega}{c^2}\int_0^{x_0} I_z^{(r)}(y)\frac{dy}{y}. \tag{30}$$

Consider a case when the resulting electrical current produced by external EEF is purely sinusoidal, $I_z^{(r)} = 0$. Than it follows from (27),(30)

$$C_1^{(r)} = 0, \quad \mathcal{E}_1^{(i)} = -C_2^{(r)}. \tag{31}$$

At $x_0 \ll 1$ we neglect logarithmic terms in (27)-(28), and terms with integrals in (27)-(30), so the details of the current distribution over the cylinder radius are of a little importance. In this approximation

$$C_2^{(r)} = \frac{\pi\omega}{c^2} I_0^{(i)}, \quad \mathcal{E}_1^{(r)} = \frac{2\omega}{c^2} \ln \frac{x_0}{2} I_0^{(i)}. \tag{32}$$

The solution for the electromagnetic field of the long wave emitted by the cylinder with the sinusoidal electric current, starting from the surface of the cylinder until the wave zone, follows from (14):

$$\begin{aligned} E_z &= \frac{\pi\omega}{c^2} I_0^{(i)} [Y_0(x) \cos\omega t - J_0(x) \sin\omega t], \\ B_\phi &= -\frac{\pi\omega}{c^2} I_0^{(i)} [Y_1(x) \sin\omega t + J_1(x) \cos\omega t]. \end{aligned} \tag{33}$$

Near the cylinder we have, with account of the expansions (21)

$$\begin{aligned} E_z &= \frac{\pi\omega}{c^2} I_0^{(i)} (\frac{2}{\pi} \ln \frac{x}{2} \cos\omega t - \sin\omega t) \\ &\approx \frac{2\omega}{c^2} I_0^{(i)} \ln \frac{x}{2} \cos\omega t = \frac{2\omega}{c^2} I_0^{(i)} \ln \frac{r\omega}{2c} \cos\omega t, \end{aligned} \tag{34}$$

$$B_\phi = -\frac{\pi\omega}{c^2} I_0^{(i)} (-\frac{2}{\pi x} \sin\omega t + \frac{x}{2} \cos\omega t) \approx \frac{2\omega}{c^2 x} I_0^{(i)} \sin\omega t = \frac{2 I_0^{(i)}}{cr} \sin\omega t. \tag{35}$$

Note, that in the near zone of the long wave the magnetic field adiabatically follows the current through the cylinder (Landau and Lifshits, 1982). The distribution of the electrical field is similar to the one of the linearly growing current in the cylinder (Bisnovatyi-Kogan, 2003). At large r we obtain from (33), (12) the expanding cylindrical wave

$$\begin{aligned} B_\phi = -E_z &= -\frac{\pi\omega}{c^2} \sqrt{\frac{2}{\pi x}} I_0^{(i)} \sin(x - \frac{\pi}{4} - \omega t) \\ &= -\frac{1}{c} \sqrt{\frac{2\pi\omega}{cr}} I_0^{(i)} \sin[\frac{\omega}{c}(r - ct) - \frac{\pi}{4}]. \end{aligned} \tag{36}$$

6. Electromagnetic energy flux from jet

Strong electromagnetic wave generated by oscillations may accelerate effectively particles at large distances from the nucleus near the jet, as well as at larger radiuses (Bisnovatyi-Kogan and Lovelace, 1995). Let us estimate the energy flux in the electromagnetic wave radiated by the jet of the length l, and radius r_0. If n_e is the electron density producing

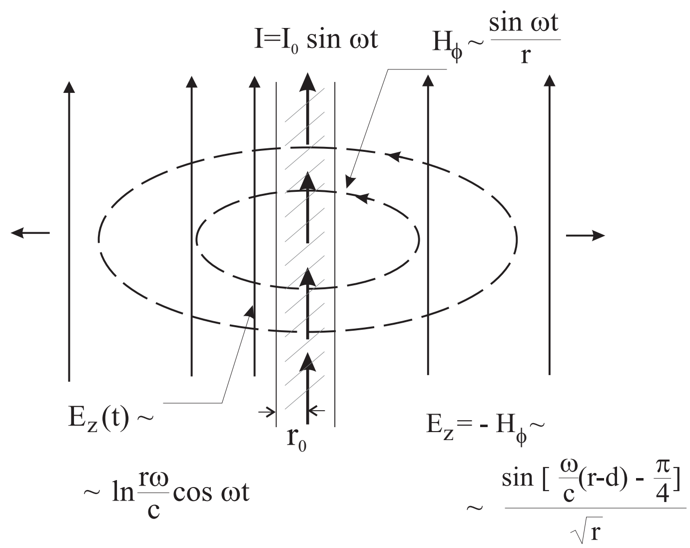

Figure 1. Magnetic and electrical fields around the infinite cylinder with the radius r_0, and low-frequency $\omega \ll c/r_0$, sinusoidal electrical current along the cylinder axis. In the near zone electrical and magnetic fields are varying in antiphase, and far from the cylinder $r \gg c/\omega$ the expanding cylindrical electromagnetic wave is formed, with $E_z = -B_\phi$.

the electrical current, with e as electron charge, than, with account of (36) the Pointing flux $P = \frac{c}{4\pi}[\mathbf{EB}]$ through the cylinder surface is

$$F = 2\pi r_0 l P = \frac{\pi l \omega}{2c^2} I_0^2. \tag{37}$$

For the amplitude of the electrical current along the cylinder radius $I_0 = \pi r_0^2 n_e c e$, we obtain the energy flux from the jet in the form:

$$F = \frac{\pi^3}{2} e^2 l r_0^4 \omega n_e^2 \approx 2 \cdot 10^{49}\ \mathrm{erg/s}\ \frac{l}{1\,\mathrm{kpc}} \left(\frac{r_0}{1\,\mathrm{pc}}\right)^4 \frac{T}{100\,\mathrm{yr}} \left(\frac{n_e}{10^{-10}\,\mathrm{cm}^{-3}}\right)^2. \tag{38}$$

Here $T = \frac{2\pi}{\omega}$ is the period of the electromagnetic wave. Part of the radiated energy is used for particle acceleration up to very large energies (Bisnovatyi-Kogan and Lovelace, 1995), and support the jet radiation at different energy regions of electromagnetic spectrum.

7. Generation of strong electromagnetic wave by proper oscillations in a separate blob.

The mechanism of shock acceleration of particles, often considered (Eilek and Hughes, 1990), is not certain, it is unlikely that shock

acceleration can give a fairly uniform brightness jet as observed in some cases. The mechanisms of magnetic field reconnection (Romanova and Lovelace, 1992) and plasma turbulence acceleration (Eilek and Hughes, 1990) are also highly uncertain. When plasma density in the surrounding medium is small, the electromagnetic wave generated by the nonpotential plasma oscillations of the confined body is emitted outside and can accelerate particles. When the emitted wave is strong enough it washed out the medium around and the density may become very small, consisting only of the accelerated particles. The solution of the whole problem of the dynamical behavior of the confined knots embedded into the large scale elongated magnetic field and producing the toroidal field can be solved by self-consistent calculations of the knot oscillations, using together the hydrodynamical and complete Maxwell equations. In order to estimate properties of a long wave radiation of by oscillating knot we solve instead the idealized problem having the analytical solution. Consider linear plasma oscillations of the infinitely long uniform cylinder. The problems of such kind have been intensively studied for plasma wave-guides (Kondratenko, 1976). The main difference in this problem is another boundary conditions which suggest vacuum state around the cylinder. When considering linear electromagnetic oscillations in the static plasma cylinder, only Maxwell equations (3) with time dependence in the form $\sim\ \exp(-i\omega t)$ are needed. The background constant field $B_z = B_0$ is adopted. Dependence if j on $\mathbf{E}$, $\mathbf{B}$ and B_0 can be obtained, using the expression for the dielectric permeability

$$\epsilon_{ij} = \delta_{ij} + \frac{4\pi i}{\omega}\left[\sigma_{ij}(e) + \sigma_{ij}(p)\right], \;\; j_i = \left[\sigma_{ij}(e) + \sigma_{ij}(p)\right] E_j \tag{39}$$

Here we consider for simplicity pure hydrogen plasma. The components for $\sigma_{ij}(\alpha)$ in the case of perfect conductivity are (Kondratenko, 1976)

$$\sigma_{11}(\alpha) = \sigma_{22}(\alpha) = \frac{i}{4\pi}\frac{\omega_{p\alpha}^2\,\omega}{\omega^2 - \omega_{B\alpha}^2}, \tag{40}$$

$$\sigma_{12}(\alpha) = -\sigma_{21}(\alpha) = -\frac{1}{4\pi}\frac{\omega_{p\alpha}^2\,\omega_{B\alpha}}{\omega^2 - \omega_{B\alpha}^2}, \quad \sigma_{33}(\alpha) = \frac{i}{4\pi}\frac{\omega_{p\alpha}^2}{\omega}$$

where $\omega_{p\alpha}$ and $\omega_{B\alpha}$ are plasma and Larmor frequencies of electrons (e) and protons (p): $\omega_{p\alpha} = \frac{4\pi n_0 e_\alpha^2}{m_\alpha}$, $\omega_{B\alpha} = \frac{e_\alpha B_0}{m_\alpha c}$. In the cylindrical coordinates $(r,\ \phi,\ z)$ with $\frac{\partial}{\partial\phi} = 0$, $\frac{\partial}{\partial z} = ik$ we have from (3),(39)

$$-ikB_\phi = \frac{4\pi}{c}\,j_r - \frac{i\omega}{c}E_r, \tag{41}$$

$$ikB_r - \frac{dB_r}{dr} = \frac{4\pi}{c} j_\phi - \frac{i\omega}{c} E_\phi, \quad \frac{1}{r}\frac{d}{dr}(rB_\phi) = \frac{4\pi}{c} j_z - \frac{i\omega}{c} E_z,$$

$$-ikE_\phi = \frac{i\omega}{c} B_r, \; ikE_r - \frac{dE_z}{dr} = \frac{i\omega}{c} B_\phi, \; \frac{1}{r}\frac{d}{dr}(rE_\phi) = \frac{i\omega}{c} B_z. \quad (42)$$

We are interested in long wave oscillations with $\omega \ll \omega_{pe}$, ω_{pp}, ω_{Be}, ω_{Bp}, so approximately

$$j_r = -\frac{i}{4\pi}\frac{\omega_{pp}^2 \omega}{\omega_{Bp}^2} E_r, \; j_\phi = -\frac{i}{4\pi}\frac{\omega_{pp}^2 \omega}{\omega_{Bp}^2} E_\phi, \; j_z = -\frac{i}{4\pi}\frac{\omega_{pe}^2}{\omega} E_z. \quad (43)$$

Substituting (43) into (41),(42) we can see that two types of long waves, corresponding to different polarizations exist independently: electric-type (E) waves with $B_z = 0$ and nonzero (E_r, E_z, B_ϕ), and magnetic-type (B) waves with $E_z = 0$ and nonzero (B_r, B_z, E_ϕ). The equations for $E-$ wave has a form

$$\frac{d}{dr}\left[\frac{1}{r}\frac{d}{dr}(rZ_E)\right] + \left(\frac{\omega_{pe}^2}{\omega^2} - 1\right)\left(\frac{k^2 c_A^2}{c^2} - \frac{\omega^2}{c^2}\right) Z_E = 0, \quad (44)$$

$$\frac{1}{r}\frac{d}{dr}\left(r\frac{dE_z}{dr}\right) + \left(\frac{\omega_{pe}^2}{\omega^2} - 1\right)\left(\frac{k^2 c_A^2}{c^2} - \frac{\omega^2}{c^2}\right) E_z = 0, \quad (45)$$

where $Z_E = E_r$ or B_ϕ, and

$$c_A^2 = c^2 \left(1 + \frac{\omega_{pp}^2}{\omega_{Bp}^2}\right)^{-1} = c^2 \left(1 + \frac{4\pi\rho c^2}{B^2}\right)^{-1} \quad (46)$$

is a speed of the Alfven waves. The equations for a $B-$ wave has a form

$$\frac{d}{dr}\left[\frac{1}{r}\frac{d}{dr}(rZ_B)\right] + \left(\frac{\omega^2}{c_A^2} - k^2\right) Z_B = 0, \quad (47)$$

$$\frac{1}{r}\frac{d}{dr}\left(r\frac{dB_z}{dr}\right) + \left(\frac{\omega^2}{c_A^2} - k^2\right) B_z = 0, \quad (48)$$

where $Z_B = B_r$ or E_ϕ, Outside the cylinder equations (44),(45), (47),(48) retain their form after substituting c everywhere instead of c_A. Solutions of these equations are given by Bessel functions:

$$Z_{B,E} = Z_{B0,E0} J_1(\kappa_{B,E} r), \; B_Z = B_{Z0} J_0(\kappa_B r), \; E_Z = E_{Z0} J_0(\kappa_E r), \quad (49)$$

where $\kappa_B^2 = \frac{\omega^2}{c_A^2} - k^2$, $\kappa_E^2 = \left(\frac{\omega_{pe}^2}{\omega^2} - 1\right)\left(\frac{k^2 c_A^2}{c^2} - \frac{\omega^2}{c^2}\right)$ inside the cylinder. Outside the same solutions (49) are valid with $\kappa^2 = \frac{\omega^2}{c^2} - k^2$, instead of

κ_E^2 or κ_B^2. Discreet values of κ_B and κ_E are determined by dispersion equation, obtained from boundary conditions on the surface of the cylinder. For $B-$ wave the components B_r and E_ϕ are continuous on the boundary and $B_z = 0$ is taken at the boundary $r = r_0 - 0$. That leads to the relations

$$\kappa_{Bn} r_0 = \lambda_{0n}, \quad \omega_n^2 = c_A^2 \left(k^2 + \frac{\lambda_{0n}^2}{r_0^2} \right). \tag{50}$$

For $E-$ wave E_z is continuous and $E_r = 0$ at $r = r_0 - 0$, so we have

$$\kappa_{Ei} r_0 = \lambda_{1i}, \quad \omega_i^2 = c_A^2 k^2 \left(1 + \frac{\lambda_{1i}^2 c^2}{r_0^2 \omega_{pe}^2} \right)^{-1}, \tag{51}$$

where λ_{0n} and λ_{1i} are zero's of the Bessel functions $J_o(x)$ and $J_1(x)$. We have then for the wave vectors outside the cylinder

$$\kappa_n^2 = \frac{c_A^2}{c^2} \frac{\lambda_{0n}^2}{r_0^2} - k^2 \left(1 - \frac{c_A^2}{c^2} \right) \ (B), \quad \kappa_i^2 = -k^2 \left(1 - \frac{c_A^2/c^2}{1 + \frac{\lambda_{1i}^2 c^2}{r_0^2 \omega_{pe}^2}} \right) \ (E). \tag{52}$$

It is clear from (46) that $\kappa_i^2 < 0$ for $E-$ waves, but for each k there is n_* such that $\kappa_n^2 > 0$ for $n > n_*$ and Alfven $B-$ wave is transforming into electromagnetic one. The vacuum values of $B_r(r_0+0)$ and $E_\phi(r_0+0)$ at the boundary are obtained from the continuity conditions, and the outer value of $B_z(r_0+0)$ is found from the last equation in (42). The jump $\Delta B_z = B_z(r_0+0)$ on the boundary determines the induced surface current i_ϕ. Very long electromagnetic wave emitted with the amplitude E_ϕ could accelerate particles up to energies $\epsilon \approx e E r_0 \approx \alpha 10^3$ ergs for $r_0 = 1$ pc and $E = \alpha \cdot 10^{-6}$ in CGS. Real energy of the accelerated particles could be much less due to radiation losses.

8. Conclusion

Similarity of the acceleration mechanisms of the particles in the pulsars and relativistic jets could be a reason of the similarity of the high energy radiation around 100 Mev which have been observed by EGRET in number of radiopulsars (Thompson et al., 1992); quasars and AGN (Thompson et al., 1993), but in no other compact objects.

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